The self-penguin contribution to $K \to 2\pi$

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Abstract

We consider the contribution to $K \to 2\pi$ decays from the non-diagonal $s \to d$ quark transition amplitude. First, we calculate the most important part of the $s \to d$ transition, the so-called self-penguin amplitude $\sim G_F \alpha_s$, including the heavy top-quark case. Second, we calculate the matrix element of the $s \to d$ transition for the physical $K \to 2\pi$ process. This part of the analysis is performed within the Chiral Quark Model where quarks are coupled to the pseudoscalar mesons.

The CP-conserving self-penguin contribution to $K \to 2\pi$ is found to be negligible. The obtained contribution to ϵ'/ϵ is sensitive to the values of the quark condensate $< \bar{q}q >$ and the constituent quark mass M. For reasonable values of these quantities we find that the self-penguin contribution to ϵ'/ϵ is 10-15 % of the gluonic penguin contribution and has the same sign. Given the large cancellation between gluonic and electroweak penguin contributions, this means that our contribution is of the same order of magnitude as ϵ'/ϵ itself.

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1. Introduction

In general, non-leptonic $\Delta S = 1$ decays are described by an effective Lagrangian at quark level [1, 2, 3]

$$\mathcal{L}(\Delta S = 1) = \sum_i C_i Q_i \ . \tag{1}$$

In such effective Lagrangians, the heavy mass scales are integrated out and their effects are contained in the Wilson coefficients C_i. The quark operators Q_i involve the three light quarks $q = u, d, s$. The coefficients will in general include short distance QCD effects calculated by means of perturbation theory and the renormalization group equations (RGE). In our notation the Wilson coefficients C_i also contain Fermis coupling constant G_F and the relevant Kobayashi-Maskawa factors $\lambda_q = V_{sq}V_{dq}^*$ for CP-conserving $(q = u)$ and CP-violating $(q = t)$ $\Delta S = 1$ transitions respectively.

Within this standard procedure, one usually omits operators containing $(i\gamma \cdot D - m_q)$, by appealing to the QCD equations of motion for quark fields [2, 4, 5]:

$$(i\gamma \cdot D - m_q) \to 0 \ , \tag{2}$$

where D_μ is the covariant derivative containing the gluon field, and m_q the appropriate quark mass. This procedure corresponds to going on-shell with external quarks in quark operators. Certainly, quarks are not on-shell in hadrons, especially not in the (would-be Goldstone) mesons π and K.

In a series of papers [6, 7, 8, 9], it has been pointed out that the effect of such operators are not zero in general, and that they in particular will contribute to processes like $K \to 2\pi$, $K \to 2\gamma$, and $B \to 2\gamma$. In order to be consistent, such genuinely off-shell effects should go to zero in the limit when some typical bound state parameter goes to zero, which is indeed found to be the case. In our case this parameter is the constituent quark mass of the light u, d, s quarks. Numerically, the non-zero off-shell effects are non-negligible in some cases.

In this paper we will study a quantity which by construction vanishes on the s- and d- quark mass-shells after regularization, namely the non-diagonal $s \to d$ self-energy transition Σ_{ds} [6, 7, 8, 10]. For off-shell quarks, however, this renormalized non-diagonal self-energy is still non-zero and could a priori be relevant for physical amplitudes, e.g $K \to 2\pi$. The unrenormalized $s \to d$ transition corresponds to an effective Lagrangian

$$\mathcal{L}_{ds}^U = \bar{d}(ai\gamma \cdot DL + bL + cR)s \ , \tag{3}$$

where a, b, c are divergent quantities given by the loop integrations, and L, R are the left- and right-handed projectors in Dirac space. In the Standard Model Lagrangian there are no direct $s \to d$ transitions. Consequently, one

demands that the renormalized self-energy vanishes on the respective mass shells of the s- and d-quarks, by adding the necessary counterterms. It should be emphasized that this defines the physical s and d quarks in the presence of weak interactions. The renormalized self-energy corresponds to an effective Lagrangian of the form [11]

$$\mathcal{L}_{ds}^R = -A\,\bar{d}(i\gamma \cdot D - m_d)(i\gamma \cdot DR + M_R R + M_L L)(i\gamma \cdot D - m_s)s \ . \quad (4)$$

In momentum space, A is a finite, slowly varying (logarithmic) function of the W-boson and heavy quark masses and the external s- and d- quark momentum (k) squared. To be explicit, A is obtained from an expansion of a in (3): $a = a_0 + k^2 A + \cdots$. The quantity a_0, which is divergent and independent of k^2, is removed by the counterterms. The quantities $M_{L,R}$ depend on the masses of the s- and d-quarks.

In the limit $m_{s,d} \to 0$, which we will work, the Lagrangian is

$$\mathcal{L}_{ds}^R(m_{s,d} \to 0) = -A\,\bar{d}(i\gamma \cdot D)^3 Ls \ . \quad (5)$$

In the pure electroweak case, there is a strong suppression of the CP-conserving amplitude owing to the GIM-mechanism [12], and one finds [1] that A is of order $G_F m_c^2/M_W^2$. In the CP-violating case this strong GIM suppression is relaxed due to a top-quark with a mass of the same order as the W-boson. Even more important, when perturbative QCD to lowest order is added, one obtains an unsuppressed charm-quark contribution $\sim G_F \alpha_s (log\,m_c)^2$ [7, 8]. A Feynman diagram for this contribution is shown in Fig. 1 and is named *self-penguin* because the gluon from a one loop penguin diagram is reabsorbed by the external s- or d-quarks.

If one applies the perturbative QCD equations of motion as in (2), we observe that $\mathcal{L}_{ds}^R \to 0$. However, we should emphasize that $\mathcal{L}_{ds}^R$ is obtained within perturbation theory. Therefore, it is not allowed to apply (2) for quarks strongly bound in π, K. An analogue within QED would be to replace (2) by

$$(\gamma \cdot D - m) \to e\gamma \cdot A^C \ , \quad (6)$$

where the perturbative photons are contained in the covariant derivative, and A_μ^C represents the binding Coulomb forces[13]. Also, in QED the renormalized electron self-energy is zero on the electron mass-shell, but still it gives an important contribution to Lamb-shift. As a gedanken experiment, if the s- and d-quarks were bound by electromagnetic forces, the contribution to $(\bar{d}s) \to 2\gamma$ from (5) would be proportional to the binding energy of the system, that is, of the order α_{em}/r_B, where r_B is the Bohr radius. Being small in QED, one expects off-shell effects to be bigger in strong interactions where bound state effects are more important. Thus, physical effects for
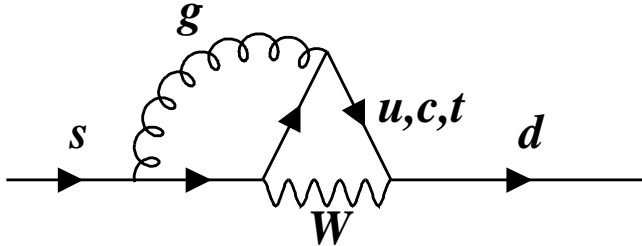
Figure 1: The self-penguin diagram. There is also a counterpart where the gluon is absorbed by the d-quark.

$K \to 2\pi$ decays from $\mathcal{L}_{ds}^R$ could be obtained, and one should explore possible consequences for the $\Delta I = 1/2$ rule and for ϵ'/ϵ.

To calculate the matrix elements of quark operators between physical hadronic states is in general a difficult task, and one normally uses various models or assumptions. In this paper we will use the Chiral Quark Model (χQM), advocated by many authors [14]. This model consists of the ordinary QCD Lagrangian and an extra term which induces meson-quark couplings that makes it possible to calculate matrix elements of quark operators in terms of quark loop diagrams. The model is thought to be applicable for quark momenta below some scale of the order Λ_χ, the chiral symmetry breaking scale. We find this model very interesting because it reproduces the chiral Lagrangian terms for strong interactions to good accuracy to order p^4. Furthermore, it has been shown to be useful in calculating weak hadronic matrix elements [15, 16].

2. The $s \to d$ self-penguin transition

The renormalized non-diagonal $s \to d$ off-diagonal self-energy has the following form in agreement with (5)

$$\mathcal{M}(s \to d) = -i\bar{d}\,\Sigma_{ds}\,s = -ik^2\,\bar{d}\,\gamma \cdot kLsA(k) \ , \tag{7}$$

where k is the four momentum of the external s- and d- quark. In the pure electroweak (EW) case we obtain for one flavour $q = u, c, t$:

$$A(k)_{EW} = 2\lambda_q \frac{g_W^2}{16\pi^2}(1 + \frac{m^2}{2M_W^2}) \int_0^1 dx \frac{x}{k^2} \ln\left[1 - \frac{x(1-x)k^2}{\overline{M_x^2}}\right] \ , \tag{8}$$

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where $\overline{M_x^2} = xM_W^2 + (1-x)m^2$, M_W is the W-boson mass and $m = m_q$ is the mass of the quark running in the loop ($q = u, c, t$), g_W is the W-coupling, and λ_q is the appropriate KM factor. The term $\sim m^2/(2M_W^2)$ is due to unphysical Higgs exchange in the Feynman gauge. For small k^2 we obtain

$$A_{EW} = \frac{G_F}{\sqrt{2}} \frac{1}{2\pi^2} \left[\lambda_u(r_u - r_c) - \lambda_t(r_c - r_t) \right] , \tag{9}$$

where the quantities r_q are given by

$$r_q = -(1 + \frac{m^2}{2M_W^2}) \int_0^1 dx \, x^2(1-x) \frac{M_W^2}{\overline{M_x^2}} , \tag{10}$$

for $q = u, c, t$. As already mentioned, the pure electroweak CP-conserving $s \to d$ transition amplitude is very suppressed due to the GIM mechanism. A priori one expects a relaxed GIM mechanism for the CP-violating case involving the heavy top quark. The top quark contribution is, in Feynman gauge, dominated by the unphysical Higgs contribution. As a curiocity, we find a "super-GIM" mechanism: It turns out that the top-quark contribution exactly cancels the charm contribution $((r_c - r_t) \to 0)$ in the limit $m_c \to 0$ and $m_t \to \infty$. Therefore the GIM-cancellation is somewhat stronger than expected in the pure electroweak CP-violating case. Numerically, we find $(r_c - r_t) = -0.167 - (-0.124) = -0.043$ for $m_t = 180$ GeV.

As shown originally by Shabalin [7], the self-penguin diagram has only a weak, logarithmic GIM suppression to order $\alpha_s G_F$. To obtain the two loop self-penguin we need the penguin ($s \to d + gluon$) loop for arbitrary external momenta. By direct calculation this is divergent, and before we insert it into the next loop to get the self-penguin, it has to be properly regularized. For this purpose we consider the Ward-identity

$$q_\nu \Gamma_{ds}^\nu = \Sigma_{ds}(k+q) - \Sigma_{ds}(k) , \tag{11}$$

in order to determine the necessary counterterms. In our case Γ_{ds}^ν corresponds to the penguin one loop (the strong coupling g_s not included). Writing

$$\Gamma_{ds}^\nu(k, q) = T^{\sigma\nu} \gamma_\sigma L , \tag{12}$$

the most general tensor compatible with the Ward identity (11) can be written

$$\begin{aligned} T^{\sigma\nu} = & \left[(k+q)^2 A(k+q) - k^2 A(k) \right] \frac{k^\sigma q^\nu}{q^2} + (k+q)^2 A(k+q) \, g^{\sigma\nu} \\ & + B_0[q^\sigma q^\nu - q^2 g^{\sigma\nu}] + B_1[q^\sigma k^\nu - q \cdot k g^{\sigma\nu}] \\ & + B_2[k^\sigma k^\nu - \frac{q \cdot k}{q^2} k^\sigma q^\nu] + B_\epsilon i\epsilon^{\sigma\nu\alpha\beta} k_\alpha q_\beta . \end{aligned} \tag{13}$$

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Comparing this with the raw result for Γ^ν_{ds}, we find the regularized expressions:

$$B_j = 2g_W^2 \int_0^1 dx \int_0^{1-x} dy f_j \left[\overline{M_y^2} - x(1-x)q^2 - y(1-y)k^2 + 2xyq \cdot k\right]^{-1}. (14)$$

For W- exchange only, the f_j's are given by $f_0 = 2x(1-x); f_1 = y(2x-1); f_2 = -2y^2$ and $f_\epsilon = -y$, respectively. There are similar expressions for the unphysical Higgs case. We obtain the following expression for the renormalized self-penguin contribution:

$$\mathcal{M}(s \to d, g_s^2) = -2ig_s^2 g_W^2 \overline{d}t^a t^a \int \frac{d^4q}{(2\pi)^4} \frac{\gamma_\nu}{q^2} S(k+q)\Gamma^\nu_{ds}(k,q)\, s\,, \qquad (15)$$

where t^a is a colour matrix. To obtain the result compatible with (5), we expand the d- (or s-) quark propagator $S(k+q)$ and Γ^ν_{ds} in powers of the external momentum k and keep in total the order $(k)^3$ only. Then we obtain (for $q = c, t$):

$$\mathcal{M}(s \to d, g_s^2) = -i\frac{g_s^2 g_W^2}{3M_W^2}\lambda_q k^2 k^\mu \overline{d}t^a t^a \gamma_\mu Ls \left(\frac{i}{16\pi^2}\right)^2 F_q\,, \qquad (16)$$

where the dimensionless quantity:

$$
\begin{aligned}
F_q \;=\;& 6M_W^2 \int_0^1 dx \int_0^{1-x} dy \left\{ \frac{[(-1/2 + 2x)y + x(1-x)]ln(\frac{\overline{M_y^2}}{x(1-x)\Lambda^2})}{[\overline{M_y^2} - x(1-x)\Lambda^2]} \right. \\
& \left. - \frac{yx}{\overline{M_y^2}} + \frac{y^2(3x-1)}{(1-x)\overline{M_y^2}} - \frac{x(1-x)}{3\overline{M_x^2}}\right\} \\
& + 3m^2 \int_0^1 dx \int_0^{1-x} dy \left\{ \frac{[(-3/2 + 2x)y + x(1-x)]ln(\frac{\overline{M_y^2}}{x(1-x)\Lambda^2})}{[\overline{M_y^2} - x(1-x)\Lambda^2]} \right. \\
& \left. - \frac{y(x+3y)}{\overline{M_y^2}} - \frac{3y^2(1-y)}{x(1-x)\overline{M_y^2}} + \frac{2y^3}{(1-x)^2\overline{M_y^2}} - \frac{x(1-x)}{3\overline{M_x^2}}\right\}\,, \qquad (17)
\end{aligned}
$$

includes the unphysical Higgs contribution $\sim 3m^2$. Since we have omitted any current mass for the s- and d-quark and expanded to third order in the external momentum k, one has to introduce a lower cut-off scale Λ. In order to match the low-energy description in terms of the χQM, its magnitude will be taken as the chiral symmetry breaking scale $\Lambda_\chi = 0.83$ GeV.

It should be noted that in the previous calculation [8], only the case $m_q^2 \ll M_W^2$ was considered. In this case only the B_0 term contributes to order G_F, while contributions from the other B_i's are suppressed by m_q^2/M_W^2

with respect to the B_0 contribution. Including the heavy top quark, all contributions have to be taken into account.

Combining the equations (16) and (17), we find the self-penguin contribution to the quantity A in (5):

$$A_{SP} = \frac{G_F}{\sqrt{2}} \frac{h}{24\pi^2} \frac{\alpha_s}{\pi} [\lambda_u(F_u - F_c) - \lambda_t(F_c - F_t)] , \qquad (18)$$

where $h_c = Tr(t^a t^a)/N_c = (N_c^2 - 1)/(2N_c)$ is a colour factor, and N_c is the number of colours. Eqs. (15)-(17) is the contribution from the diagram in Fig. 1 (including the unphysical Higgs contribution). Taking into account the counterpart where the gluon is absorbed by the d-quark, one gain a factor two which is included in (18). Numerically, we find for $m = m_c = 1.4$ GeV and $\Lambda = 0.83$ GeV:

$$F_c \simeq 49.0 . \qquad (19)$$

Similarly, for $m = m_t = 180$ GeV and $\Lambda = 0.83$ GeV, we obtain

$$F_t \simeq -10.7 . \qquad (20)$$

For the c-quark case one obtains the analytical result within the leading logarithmic approximation

$$F_c = \frac{1}{2}[\ln(\frac{M_W^2}{\Lambda^2})]^2 - \frac{1}{2}[\ln(\frac{m_c^2}{\Lambda^2})]^2 + (\frac{5}{6} - \frac{4}{6})\ln(\frac{M_W^2}{m_c^2}) - \frac{5}{6}\ln(\frac{m_c^2}{\Lambda^2}) . \qquad (21)$$

The leading part of this result was obtained in [8] for the charm quark case, while the general result (17) is new. In (21) the term $\sim -4/6$ is coming from $-B_0 q^2 g^{\sigma\nu}$ in (13)and the rest from $B_0 q^\sigma q^\nu$. As shown in [17, 18], leading logarithmic QCD corrections can be taken into account by folding the well known Wilson coefficients $C_\pm$ [1, 2] into an integral over virtual momenta p. Then our result takes the form

$$\frac{\alpha_s}{\pi} F_c \rightarrow \tilde{F}_c = \int_{m_c^2}^{M_W^2} \frac{dp^2}{p^2} \rho(p^2) \left\{ \ln(\frac{p^2}{\Lambda^2}) + (\frac{5}{6} - \frac{4}{6}) \right\} - \frac{5}{6}\rho(m_c^2) \ln(\frac{m_c^2}{\Lambda^2}) , \qquad (22)$$

where $\rho(p^2) = [C_+(p^2) + C_-(p^2)]\alpha_s(p^2)/(2\pi)$. Because the operator in (5) has zero anomalous dimension, there is no extra factor attached to $\rho(p^2)$. The linear combination $(C_+ + C_-)/2$ is the colour octet combination of the coefficients which occurs because the gluon within the self-penguin is a colour octet. The $C_\pm$'s are given by ratios of α_s at different scales to a power determined by the anomalous dimension of the corresponding four quark operators $Q_\pm$. Thus the integral in (22) can be performed analytically, and the result obtained in terms of α_s at the scales $\Lambda, m_c, m_b,$ and M_W respectively.

Numerically, we find $\tilde{F}_c \simeq 3.25$. For the top quark contribution, we expect essentially no QCD corrections except for the running of α_s, and we take

$$\tilde{F}_t = \frac{\alpha_s(m_t^2)}{\pi} F_t \simeq -0.37 \; , \tag{23}$$

approximately the same as the pure EW contribution. The total A to be used in (5) is $A_{SP} + A_{EW}$. (There are other contributions of order $\alpha_s G_F$, but with a very small numerical coefficient. Note also that A_{EW} is dominated by loop momenta of order M_W where $C_\pm \simeq 1$, and there will be no further QCD corrections in the leading logarithmic approximation).

3. The Chiral Quark Model

In the Chiral Quark Model (χQM) [14], chiral-symmetry breaking is taken into account by adding an extra term $\mathcal{L}_\chi$ to ordinary QCD:

$$\mathcal{L}_\chi = -M \left(\overline{q_L} \Sigma^\dagger q_R + \overline{q_R} \Sigma q_L \right) \; , \tag{24}$$

where $\bar{q} = (\bar{u}, \bar{d}, \bar{s})$. The constant M in (24) is interpreted as the *constituent* quark mass, expected to be of order 200-300 MeV. The quantity Σ contains the Goldstone-octet fields π_a:

$$\Sigma = exp(i \sum_a \lambda_a \pi_a / f) \; , \tag{25}$$

where the λ_a's are the Gell-Mann matrices, and $f = f_\pi = 93$ MeV is the pion decay constant. The term $\mathcal{L}_\chi$ introduces a meson-quark coupling. This means that the quarks can be integrated out and the coefficients of the various terms in the chiral Lagrangian are calculable [14].

The χQM should be applied for momenta lower than the scale of chiral-symmetry breaking, $\Lambda_\chi = 2\pi f_\pi \sqrt{6/N_c}$, where $N_c = 3$ is the number of colours. Divergent integrals are regularized, for instance with dimensional regularization [16]. An alternative is to use some ultraviolet cut-off Λ_U of the order Λ_χ. Because f_π, entering the meson- quark coupling $\sim M\gamma_5/f$, is also given by a quark loop diagram for $\pi \to W$(virtual), one has:

$$f_\pi^{(0)} = \frac{N_c M^2}{4\pi^2 f} [\; \hat{f}_\pi \; + \frac{\pi^2}{6 N_c M^4} < \frac{\alpha_s}{\pi} G^2 > + \cdots] \; , \tag{26}$$

where $< \frac{\alpha_s}{\pi} G^2 >$ is the two gluon condensate, and the dots indicate higher gluon condensates. In the end both f and $f_\pi^{(0)}$ will, in the limit $m_{u,d} \to 0$, be identified by f_π, but at intermediate stages one might need to distinguish them for technical reasons. The dimensionless quantity $\hat{f}_\pi$ has the leading behaviour $\sim \ln(\Lambda_U^2/M^2)$ in a cut-off type regularization and $\sim \Gamma(\epsilon) \sim 1/\epsilon$ within dimensional regularization ($D = 4 - 2\epsilon$ being the dimension of space).
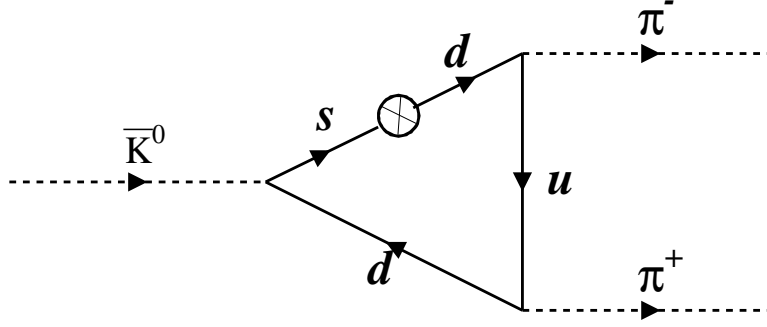
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Figure 2: Typical diagram for non-diagonal $s \to d$ self-energy induced $K \to 2\pi$ transition.

The quark condensate can be written as

$$< \bar{q}q > = \frac{N_c M}{4\pi^2}C_q - \frac{1}{12M} < \frac{\alpha_s}{\pi}G^2 > + \cdots , \qquad (27)$$

The quantity C_q depends on the regularization prescription. Using a cut-off Λ, one finds $C_q = -\Lambda^2 + M^2 \ln(\Lambda_U^2/M^2) + \cdots$, and within dimensional regularization, $C_q = -M^2\Gamma(-1+\epsilon) + \cdots$.

4. The $K \to 2\pi$ amplitude from $s \to d$.

Within the χQM, the $K \to 2\pi$ amplitude obtained from the $s \to d$ transition in (5) is determined by diagrams like the one in Fig. 2. We find the following result:

$$\mathcal{M}(\overline{K^0} \to \pi^+\pi^-)_\Sigma = \frac{\sqrt{2}}{4f_\pi^3}(m_K^2 - m_\pi^2)AM^2 \left[\frac{< \bar{q}q >}{M} + 2f_\pi^2 \left(1 - \frac{3M^2}{2\Lambda_\chi^2}\right)\right] , \qquad (28)$$

where A is the quantity in (5). For completeness, we have also calculated the virtual $K \to vacuum$ (which is zero) and $K \to \pi$ transition amplitudes in addition to the physical $K \to 2\pi$ amplitude. These calculations show that $\mathcal{L}_{ds}^R$ in (5) gives an octet contribution that can be simply added to the Q_6 contribution up to $\mathcal{O}(p^2)$. To order $\mathcal{O}(p^4)$ we expect to obtain different contributions from Q_6 and $\mathcal{L}_{ds}^R$ due to the extra derivatives in $\mathcal{L}_{ds}^R$.

We want to compare the result (28) with the corresponding amplitude for the penguin operator Q_6 :

$$\mathcal{M}(\overline{K^0} \to \pi^+\pi^-)_{Q_6} = \frac{2\sqrt{2}}{f_\pi}(m_K^2 - m_\pi^2) C_6 \frac{< \bar{q}q >}{M} \left[1 - 6\frac{M^2}{\Lambda_\chi^2}\right] . \qquad (29)$$

We observe that the ratio between the amplitudes in (28) and (29) goes to zero in the limit where the constituent quark mass $M \to 0$, as expected.

Numerically, the ratio between the self-penguin and the Q_6 contribution is less than 5 % in the CP-conserving case. For ϵ'/ϵ, the situation is different because there is a large cancellation between the gluonic and electroweak penguin contributions[19, 20, 16]. In the operator language, it is a large cancellation between the Q_6 and the Q_8 contributions. We have compared the $s \to d$ self-energy contribution with the standard Q_6 contribution in the CP-violating case. Using, in the notation of [20], $C_i = -(G_F/\sqrt{2})(\lambda_u z_i - \lambda_t y_i)$, with $y_6 = -0.137$ at the renormalization scale $\mu = \Lambda = \Lambda_\chi$; $M \simeq 220$ MeV and $< \bar{q}q >^{1/3} \simeq -220$ MeV [16], we find that in the CP-violating case, the $s \to d$ contribution is 10 to 15 % of the Q_6 contribution and has the same sign. This means that after the partial cancellation between the Q_6 and the Q_8 contributions has been taken into account, the $s \to d$ contribution is of the order ϵ'/ϵ itself.

5. Discussion

We have considered the $s \to d$ self-energy contribution to $K \to 2\pi$ within the chiral quark model. Within this model matrix elements between the light mesons (π, K) of quark operators may be calculated in terms of quark loops. Therefore this model provides a means to calculate effects of off-shell quarks.

In the literature, terms which vanish by using the perturbative equations of motion are omitted from the analysis from the very beginning[2, 4, 5]. But to neglect such effects would, within QED, be analogous to discarding Lamb-shift. Moreover, such effects are larger in QCD where the bound state effects are stronger. Still, also in strong interactions such effects are relatively small, being proportional to $M^2/\Lambda_\chi^2 \simeq 0.07$ and vanishing in the free quark limit $M \to 0$. They might, however, become non-negligible in cases where there are cancellations between the potentially largest terms, as shown in this paper for ϵ'/ϵ , and for $K \to 2\gamma$ [9] where the potentially large pole contributions cancel when the Gell-Mann-Okubo mass formula is used.

Numerically, we have found a small effect for the CP-conserving $K \to 2\pi$ amplitude. However, for the CP-violating $K \to 2\pi$ amplitude and thereby for ϵ'/ϵ the effect is non-negligible, being of the order ϵ'/ϵ itself, because of the strong cancellation between the Q_6 and Q_8 contributions. It would of course have been an advantage if $\mathcal{L}_{ds}^R$ were included in the basis of a complete next to leading QCD corrections [20], but this would have required a three loop calculation. The self-penguin contribution drives ϵ'/ϵ more to the positive side. However, due to the uncertain value of the quark condensate it is still difficult to make a very precise prediction of ϵ'/ϵ within the chiral quark model[16].

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